Princeton University
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Anomalous Heating of Plasmas
by Laser Irradiation

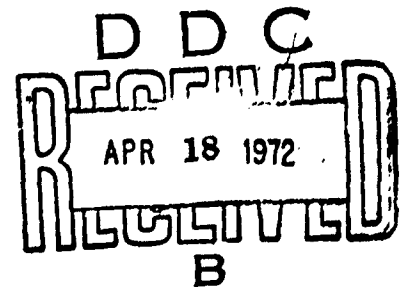by

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MATT-817   December 1970

F44620 - 70 - C-00 33

This paper was presented at the International Conference
on Laser-Produced Plasmas, Moscow, USSR, November
1970

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APR 18 1972

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AEC RESEARCH AND DEVELOPMENT REPORT

DOCUMENT CONTROL DATA - R & D

(Security classification of title, body of abstract and indexing annotation must be entered when the overall report is classified)

1. ORIGINATING ACTIVITY (Corporate author)	2a. REPORT SECURITY CLASSIFICATION
The Trustees of Princeton University Princeton, New Jersey 08540	Unclassified
	2b. GROUP

3. REPORT TITLE

Anomalous Heating of Plasmas by Laser Irradiation

4. DESCRIPTIVE NOTES *(Type of report and inclusive dates)*
Scientific   Interim

5. AUTHOR(S) *(First name, middle initial, last name)*

P. Kaw, J. Dawson, W. Kruer, C. Oberman, and E. Valeo

6. REPORT DATE	7a. TOTAL NO. OF PAGES	7b. NO. OF REFS
December 1970	26	23

8a. CONTRACT OR GRANT NO	9a. ORIGINATOR'S REPORT NUMBER(S)
b. PROJECT NO. F44620-70-C-0033	
9767-03	
c. 61445014	9b. OTHER REPORT NO(S) (Any other numbers that may be assigned this report)
d. 681301	AFOSR - TR - 72 - 0 6 5 9

10. DISTRIBUTION STATEMENT

Approved for public release; distribution unlimited

11. SUPPLEMENTARY NOTES PROCEEDINGS Int.Conf. on Laser-Produced Plasmas Moscow,USSR, November 1970	12. SPONSORING MILITARY ACTIVITY AFOSR (NPP) 1400 Wilson Boulevard Arlington, Virginia 22209

13. ABSTRACT

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Recently the problem of anomalous high-frequency resistivity and heating of plasmas has attracted considerable attention. There are two basic reasons for this interest:

(1) With the advent of high-intensity laser technology, one can readily achieve high-frequency field strengths such that the oscillating directed component of electron velocity becomes comparable to or greater than the electron thermal velocity; i.e.,

$$\frac{eE}{m \omega v_{the}} \gtrsim 1 \quad . \tag{1}$$

Qualitatively, one expects that the conventional collisional absorption processes might be strongly modified when such large field strengths are incident on a plasma. Furthermore, as will be shown below, field strengths that are orders of magnitude lower than those required by Eq. (1) are already large enough to give strong anomalous absorption effects.

(2) The conventional collisional absorption is weak at relatively high temperatures and becomes increasingly inefficient as the temperature rises. Thus the classical absorption length for frequencies $\omega \approx \omega_p$ is given by

$$L \simeq 5 \times 10^{18} \frac{T^{3/2}}{n} \quad \text{microns} \quad , \tag{2}$$

where T is the electron temperature in eV and n is the density in cm^{-3}.

For a typical ruby-laser-produced plasma with a density of 10^{21} cm^{-3} and a temperature of 5 kV the classical absorption length is 1500 microns, whereas the typical plasma radius is only about 100 microns. Under these conditions most of the laser beam energy would either be perfectly reflected or perfectly transmitted, depending on whether ω is less than or greater than the peak plasma frequency.

The above two factors strongly urge one towards an investigation of the anomalous high-frequency resistivity of a plasma. In order to really appreciate the anomalous absorption processes, however, it is important that one clearly understands the essential results of the linear theory of high-frequency resistivity.

Several authors have discussed this in considerable detail. Here we shall present some results obtained by Dawson and Oberman[1,2] in 1962. Similar results have been obtained by Perel' and Eliasberg[3] and Silin.[4] Dawson and Oberman[1,2] used a simple model for computing the high-frequency resistivity of a plasma, in which the electrons were treated as a Vlasov fluid and the ions as infinitely massive discrete scattering centers. The results of this model were later shown to be completely identical to those obtained from a rigorous theory using the BBGKY hierarchy.[5] The main results are summarized in Fig. 1, which shows a plot of high-frequency resistivity versus frequency. For low frequencies, $\nu_{ei} < \omega < \omega_p$, the resistivity is found to be a constant. For frequencies greater than about 1.4 ω_p, there is a drop in the resistivity that can be attributed to a reduction in the maximum impact parameter from (v_{the}/ω_p) to (v_{the}/ω). Near the

plasma frequency, the resistivity shows an interesting bump that can be interpreted as arising due to the excitation of longitudinal plasma waves because of an interaction between the oscillatory electron current and the ion correlations. Under equilibrium conditions this bump is small—about 1 per cent or so. However, Dawson and Oberman[2] found the interesting result that in the presence of strong ion correlations this bump can be greatly enhanced—typically by many orders of magnitude. They estimated that if there is a spectrum of ion waves peaked around wave-vector $\bar{k}$ with a width Δk, the enhancement of the resistivity is by a factor

$$\beta \simeq \frac{\delta^2}{4\pi} \frac{\bar{k}}{\Delta k} (n_i \bar{\lambda}^3) \quad , \tag{3}$$

where δ^2 is the ratio of mean square ion density fluctuation to mean ion density squared, and $(n_i \bar{\lambda}^3)$ is the number of ions in a cubic wavelength. This last factor can be very large (easily greater than 10^6), and thus β can be a large factor even with moderate values of δ^2.

Thus, even the linear theory tells us that if we produce strong ion correlations in a plasma we can enhance the high-frequency resistivity around the plasma frequency by many orders of magnitude. Let us understand the physics of this enhancement more clearly. Consider the electrons as a charged fluid and the discrete ions as potential spikes in this fluid. When we apply a uniform oscillating field, the electron fluid begins to oscillate and the stationary potential spikes jiggle it at various points. If the frequency of oscillation is close to ω_p, the spike jiggles are at the right frequency to excite the natural modes of the electron fluid. Thus a part of the electron

directed energy is lost to these plasma waves, which are then damped away.
Now, if the ion spikes are randomly distributed, then the waves they produce
will be randomly interfering with each other and not much energy can be
transferred to the waves. On the other hand, when the ions are strongly
correlated in position and/or time, they produce plasma waves which con-
structively interfere and a large amount of oscillation energy can be given
to the waves. In this case, therefore, a large part of the electron directed
energy is lost, and this shows itself as a large enhanced resistivity. Another
familiar way of stating the same fact is that there is a resonant mode coupling
of the external uniform field (k $\simeq$ 0 in the dipole approximation) to the
finite-k plasma waves through interaction with the finite-k ion wave modes.
Thus, the physical mechanism of enhanced energy absorption is the conversion
of the electromagnetic wave into longitudinal plasma waves; the ion density
fluctuations are crucial to this coupling process.

To use this effect in practice one immediately asks, "What is a good way
to produce strong ion correlations in a plasma?" The answer is provided by
the laser itself; let us see how this is done. In recent years a number of
authors—notably Silin and his co-workei [6] Dubois and Goldman,[7] Nishikawa,[8]
Sanmartin,[9] and others[10]—have investigated the excitation of low-frequency
instabilities in a plasma due to the presence of strong oscillating electric
fields, especially near the plasma frequency. One can give a simple descrip-
tion of these instabilities, known generally as parametric instabilities, in
terms of fluid equations.

When a large, long-wavelength, oscillating electric field is applied to a plasma, the equilibrium motion of electrons is given by

$$\dot{v}_o = - (eE_o/m) \cos \omega_o t \quad ; \tag{4}$$

i.e., the electron fluid is oscillating back and forth with a large velocity. Since the ions are much heavier, their equilibrium velocities are negligibly small and may be ignored. Thus our equilibrium consists of electrons streaming past ions with an oscillating velocity. For investigating the stability of this equilibrium, we can write down the equations for linearized perturbations. When the amplitude E_o is not too large, these equations may be put in the form

$$\ddot{n}_k^f + \nu_e \dot{n}_k^f + \omega_R^2 n_k^f = \frac{ike}{m} E_o(t) N_k \quad , \tag{5a}$$

$$\ddot{N}_k + \nu_i \dot{N}_k + \omega_{ia}^2 N_k = - \frac{ike}{M} \langle E_o(t) n_k^f \rangle \quad , \tag{5b}$$

$$n_k^s \simeq N_k \quad , \tag{5c}$$

where $\omega_R^2 = \omega_{pe}^2 + 3k^2 v_{the}^2$, $\omega_{ia}^2 = k^2 (KT_e/M)$, ν_e and ν_i are respectively the damping rates of high-frequency electron and low-frequency ion modes and, following Nishikawa,[8] we have separated the electron density into its slow (s) and fast (f) components. For the ions the slow motion is the dominant one, i.e., $N_k \simeq N_k^s$. In the absence of E_o, these equations represent the two natural modes of oscillation of a uniform plasma; viz., the damped plasma and ion acoustic waves. The electric field introduces a coupling between the two. Specifically, E_o interacts with the ion acoustic wave and drives the plasma wave; similarly, E_o interacts with the plasma wave and drives the

ion acoustic wave. Thus, these two natural modes pump each other in the plasma at the expense of the external field. When the rate of pumping exceeds the natural damping of the modes, the modes become unstable.

Mathematically, one can show this by letting $N_k \propto \exp(-i\omega t)$ and substituting for N_k in Eq. (5a); this leads to

$$\ddot{n}_k^f + \nu_e \dot{n}_k^f + \left[\omega_R^2 - \frac{k^2 e^2 E_o^2}{2mM(\omega_{ia}^2 - \omega^2 - i\omega\nu_i)} (1 + \cos 2\omega_o t) \right] n_k^f = 0 \quad . \quad (6)$$

Everybody will recognize this as the well-known Mathieu equation which is characteristic of parametrically driven oscillators. Instabilities are known to result when

$$2\omega_o \simeq 2\frac{\omega_{nat}}{n} \simeq \frac{2\omega_R}{n} \quad ,$$

and are strongest for $n = 1$. Interestingly enough, it turns out that two different kinds of instabilities are possible.

(1) $\omega_o > \omega_R$. In this case $\omega \simeq \Omega + i\gamma$, where, for weak fields, the real part $\Omega \simeq \omega_{ia}$; thus in this case an ion acoustic mode and an electron plasma wave both grow in the system.

(2) $\omega_o < \omega_R : \omega = i\gamma$. In this case an electron plasma wave and a purely growing ion mode are both excited in the plasma.

Both of these instabilities have definite minimum thresholds for excitation. As an example, for case (2) the minimum threshold field is given by

$$\frac{eE_o}{m\omega_o v_{the}} \simeq 2 \left(1 + \frac{T_i}{T_e}\right)^{1/2} \left(\frac{\nu_e}{\omega_p}\right)^{1/2} \quad . \quad (7)$$

Note that $(\nu_e/\omega_p)^{1/2}$ can be a very small quantity, and so the threshold field is considerably less than that given by Eq. (2). The threshold for case (1) is often smaller than that given by Eq. (7).

Although both instabilities come from the same set of equations, the purely growing one has some interesting similarities to the usual d.c. electron-ion streaming instability, and this has led some of us[9,10] to call it the "oscillating two-stream instability." The main feature of similarity with the d.c. case is that the expression for the growth rate γ, the threshold velocity v_E, and the maximum k excited (k_{max}) all tend to go over into the d.c. when one takes the limit $\omega_o \rightarrow 0$. Sanmartin[9] has given a good discussion of this.

On the basis of the above description, we expect that an intense oscillating electric field near the electron plasma frequency should drive low-frequency ion instabilities, ultimately producing strong ion correlations and hence an enhanced resistivity around the plasma frequency; this should then lead to anomalous heating of the plasma. In order to check these physical ideas we carried out a computer simulation experiment.[11]

The simulation experiment was carried out on a one-dimensional plasma with 10^4 ions and 10^4 electrons, and consisted in following the motions of ions and electrons in their own self-consistent fields and an external driver field $E = E_o \cos \omega_o t$, with $\omega_o = \omega_p$. The electron-to-ion temperature ratio was 30 and the mass ratio was 100. The choice of the frequency was motivated by

the fact that in this case ω_o is less than all $(\omega_p^2 + 3k^2 v_{the}^2)^{1/2}$ in the system, and so the only possible instability is the oscillating two-stream instability. Recent simulations with other values of ω_o close to ω_p have led to substantially the same results. We followed this plasma for several hundred plasma periods and obtained the results shown in Fig. 2. The top trace is a plot of the total wave energy in the self-consistent fields vs t, and the bottom trace is a plot of the plasma kinetic energy in random motions or its effective temperature vs t. Note that the wave energy starts exponentiating and saturates at a level two orders of magnitude higher than its initial value. The lower trace shows that initially the plasma is heating slowly, essentially by normal joule heating; however, once the ion fluctuations become sufficiently large, the plasma begins to heat very efficiently. Effective anomalous collision frequencies as high as

$$\frac{\nu^*}{\omega_p} \simeq 0.2$$

were observed. Our physical model also makes a definite prediction about the electron velocity distribution. In our mechanism, the heating takes place via the Landau damping of longitudinal plasma waves generated by the instability or the ones to which the unstable plasma waves couple by mode-coupling processes. Most of the energy should therefore be found in the tails of the electron velocity distribution. Figure 3 shows the electron velocity distribution after 250 plasma periods. Notice the large tails on the distribution consistent with the above picture. Finally we also looked at the variation of γ_{max} and k_{max} with E_o, and found the results to be in reasonable agreement with the theory given by Nishikawa.[8]

Qualitatively, our expectations have thus been vindicated. As a matter

of fact, we can even claim some quantitative agreement; thus, if one uses the

magnitude of saturated ion density fluctuations observed in the experiment

and plugs it into Eq. (3), one predicts an enhancement factor for resistivity

which agrees with the measured value within a factor of 2. However, the

basic question still remains. What is the dominant saturation mechanism

for the instability and what is the final level of ion density fluctuations? These

questions can only be answered by a fully nonlinear theory, and we have not

completed it yet. However, the nonlinear processes that we feel are dominant

(which we are investigating in detail) are the following:

(1) Mode Coupling Effects: As the electron and ion modes grow they

interact with each other via resonant mode coupling, and send energy out of

the unstable part of k space. This is a process similar to the one which

gives the anomalous damping of the external electromagnetic wave.

Another nonlinear process (which could be important) is nonlinear

Landau damping, where the beat between two high-frequency waves resonates

with the particles and leads to absorption:

$$\omega_1 - \omega_2 \simeq k v_{the} \text{ or } k v_{thi} \quad .$$

(2) Modification of the Background: This could act in many ways to saturate

the instability. Recent calculations by our group[12] have shown that nonresonant

diffusion of particles (which pulls the tails of the electron velocity distribution)

and the distinct phase correlations between unstable waves of the same k

traveling in opposite directions, leads to a modification of the natural
resonant frequency of the plasma. This then leads to a "detuning" effect
whereby the external field frequency is no longer in resonance, which finally
shuts off the instability.

Other background effects that can become important are:

(i) modifications of the zeroth-order particle orbits, when the electric
field of excited waves $E_W \simeq E_o$;

(ii) the effect of field depletion arising because of absorption of energy
from the pump—such an effect, which should be important for saturation
when E_o is only slightly greater than the threshold field, is being investi-
gated in detail by Dubois and Goldman.[13]

(3) Trapping Effects: For the strongly driven cases we feel that the
saturation of the wave amplitude is governed by the trapping of electrons in
the large-amplitude, high-frequency plasma waves. Once the wave amplitude
becomes large enough to trap a significant number of particles, it is very
difficult to increase its amplitude any further since it constantly keeps on
delivering the energy to the particles in the distribution. Phase space plots
for the electrons in some simulation experiments give clear evidence of
trapping phenomena.

In the absence of a detailed quantitative theory incorporating some or all
of the above effects, we have developed a semiquantitative model which gives
us reasonably good agreement with the simulations. We assume that as the
ion fluctuations grow they lead to an enhanced damping, not only of the
external electromagnetic wave but also of the growing plasma waves. This
can be included by replacing the electron collision frequency ν_e by $\nu^{*'}$ in

the expression for the threshold field. Thus, as the ion waves grow $\nu^{*'}$ grows, and so does the threshold field for keeping the instability going. A stage will come when $\nu^{*'}$ becomes so large that E_o no longer exceeds the threshold for instability. At this point the instability shuts itself off. For the oscillating two-stream instability, one therefore finds

$$\frac{\nu^{*'}}{\omega_p} = (eE_o/2m\omega_o v_{th})^2 \qquad . \qquad (8)$$

We now make the plausible assumption that the ν^* for the external electromagnetic wave is of the same order as the $\nu^{*'}$ for growing plasma waves. In other words,

$$\frac{\nu^*}{\nu^{*'}} \simeq O(1) \quad .$$

If we compare the results obtained from this relation with the simulation results shown in Fig. 4, they are seen to be in good agreement. For large fields there is a strong deviation from our results; we expected this since in this range trapping effects become important. It is of interest to note that for very large fields

$$\frac{\nu^*}{\omega_p} \simeq O[(\frac{m}{M})^{1/2}] \quad .$$

This seems to be a general result, and also has been verified by simulations with larger mass ratios.

We have also carried out simulation experiments similar to the above in two dimensions. The plasma heats both along the electric vector and across it. The heating rate along the electric vector is substantially the

same as in the one-dimensional case. The heating across the electric vector is at a slower rate.

We have recently carried out simulations with transverse fields of finite wavelength at $\omega_o = 2\omega_p$.[14] This also leads to an instability, as was predicted by Jackson[15] and Goldman,[16] and gives an enhancement of the bump in the resistivity at $\omega \approx 2\omega_p$ arising because of e - e interactions. The entire effect, however, is much weaker than the one discussed above. It is of interest in laser-produced plasmas, since it offers a possibility of heating underdense plasmas in an efficient manner.

We should like to mention here the work being done by Silin and his co-workers[6] on the anomalous heating by VHF fields ($\omega \gg \omega_p$). In this range the only parametric instabilities possible are some kinetic instabilities which we missed in our earlier fluid description. These instabilities are considerably weaker than the ones discussed by us, but do offer the interesting possibility of anomalously heating even an underdense plasma, which could be very useful. Simulations of these effects by Byers[17] have verified Silin's linear theory of these instabilities.

We shall devote the last part of this paper to the growing experimental evidence for the anomalous absorption effect. Most of this evidence is based on experiments conducted with microwaves and other longer-wavelength radio waves. Presumably this is because there are few experiments in which the absorption of lasers has been studied in a controlled manner.

The first piece of evidence is an experiment carried out by Gekker and Sizukhin.[18] A schematic of the experimental arrangement used by these authors is shown in Fig. 5. Powerful traveling electromagnetic waves in the H_{11} mode were launched from one end of a cylindrical waveguide, from the other end of which a plasma stream was injected inwards. The plasma density profile was typically of the form shown, and the experiment consisted in measuring the reflection coefficient as a function of the incident microwave power. Gekker and Sizukhin found the interesting result that the reflection coefficient was about 100 per cent for weak incident waves but dropped down to a low value of about 10 per cent or so very rapidly, when the incident power exceeded a certain critical value. One possible interpretation of this experiment can be given, as follows: the electromagnetic wave travels almost absorption-free until it reaches the region where the incident frequency matches the local plasma frequency. Furthermore, if it is sufficiently strong it excites an instability in this region that leads to an anomalous collision frequency ν^*, which can be estimated by using our semiquantitative model based on simulations. When ν^* becomes sufficiently large, there is a significant absorption of the incident microwave which leads to a reduction in the reflection coefficient. Using the WKB approximation for the reflection from an inhomogeneous plasma, the expression for ν^* derived above, we have obtained[19] a theoretical curve for the variation of $|R|^2$ with E_o which at low field strength is in reasonable agreement with the experiment (Fig. 6). We realize that the use of the WKB approximation and a particular form

for ν^* can be questioned; however, the purpose of this exercise was to

bring out the physical features of the anomalous process and to stimulate

further experiments.

Following the lead of Gekker and Sizukhin, Eubank[20] of Princeton has

recently carried out the following experiment (Fig. 7). Powerful micro-

waves are incident radially on a plasma stream traveling along a magnetic

field. The electric vector is polarized in the direction of the magnetic field.

The plasma is diagnosed by using probes that detect ion waves, and its

temperature is determined by diamagnetic loop measurement. Eubank has

found that when the incident power exceeds a critical value, ion waves are

excited and a significant heating of the plasma takes place. When no ion

waves are excited, there is no heating; this seems to be the first direct

evidence of the association between the excitation of ion waves and the

anomalous heating.

Some more experimental evidence[21] recently came from an entirely

different quarter; viz., from experiments on artificial heating of the

ionosphere at Boulder, Colorado (Fig. 8). These experiments consisted in

irradiating the ionospheric F layer with radio waves (having frequency less

than the critical frequency) from a powerful transmitter and observing the

influence on the propagation characteristics of a low-power exploring wave

of slightly different frequency reflected from the same region of the ionosphere.

It was expected that the powerful radio waves would heat the ionosphere and

hence decrease the electron-ion collision frequency, thus leading to a

decreased attenuation of the exploring wave. Experimentally it was found, however, that a stronger extinction of the ordinary mode of the exploring wave took place while the transmitter was on. Recent estimates by our group[22] have shown that the radio wave power in the reflection region was sufficiently strong to excite parametric instabilities producing strong ion correlations. This then leads to an enhanced high-frequency resistivity around the electron plasma frequency which is consistent with an enhanced attenuation of the exploring ordinary wave. Furthermore, the observation that the exploring extraordinary mode was relatively unaffected is consistent with the fact that its propagation path did not lie in the unstable region.

A final experiment is due to Dreicer and his co-workers,[23] who presented some preliminary results at an APS meeting recently (Fig. 9). A powerful microwave field was set up in a cavity resonator of the type shown. A thin stream of overdense plasma ($\omega_p > \omega$) was passed through the cavity; the thickness of the stream was less than the skin depth, so that the field penetrated the plasma. The experiment consisted in measuring the Q of the cavity as a function of microwave power. The following interesting result was obtained. For low powers the Q increased with power, which is consistent with the classical picture of decreasing dissipation due to electron heating; however, when the power exceeded a certain critical value the Q decreased very rapidly, showing strong anomalous dissipation. The threshold power was found to be higher than the minimum threshold for the excitation of parametric instabilities. No further details of the experiment are available at the moment.

In conclusion, one can say that the existence of an anomalous high-frequency resistivity and heating has been demonstrated without doubt by recent computer simulations and laboratory experiments with microwaves. The qualitative understanding of the anomalous resistivity is quite good; one can even make some semiquantitative estimates. A detailed nonlinear theory is being developed and should enable making definite quantitative predictions. From the point of view of controlled thermonuclear fusion research involving laser-produced plasmas, there are many unanswered questions, some of which I shall pose here:

(1) Why have not the anomalous absorption effects been observed in laser plasma experiments so far? Is it that the laser powers are not high enough, or does the presence of electron density and temperature gradients have a strong influence on these instabilities? Experimentally, one may learn something by comparing the laser plasma experiments with the microwave experiments and investigating the important differences between the two.

(2) Can this large energy, which is primarily given to electrons, be transferred to the ions within nanoseconds (which is the typical time the plasma takes to expand considerably)? In particular, are ther available any anomalous mechanisms of energy relaxation between electrons and ions?

(3) Will this heating be highly localized near the $\omega \approx \omega_p$ layer, or is there strong normal or anomalous thermal conductivity operative so as to make the temperature uniform?

One can also ask the interesting question: Does this anomalous high-frequency resistivity raise serious doubts about the feasibility of conventional RF confinement schemes, where one makes calculations assuming normal ohmic dissipation? A simple calculation based on the pressure balance argument $B^2/4\pi = nT$ and a skin penetration depth of c/ω_p leads to an oscillating current

$$j = n e v_D \approx n e v_T \quad ,$$

which is sufficiently large to excite these instabilities. This means that the dissipation might be anomalous and the power requirements too high to be of sufficient interest.

ACKNOWLEDGMENTS

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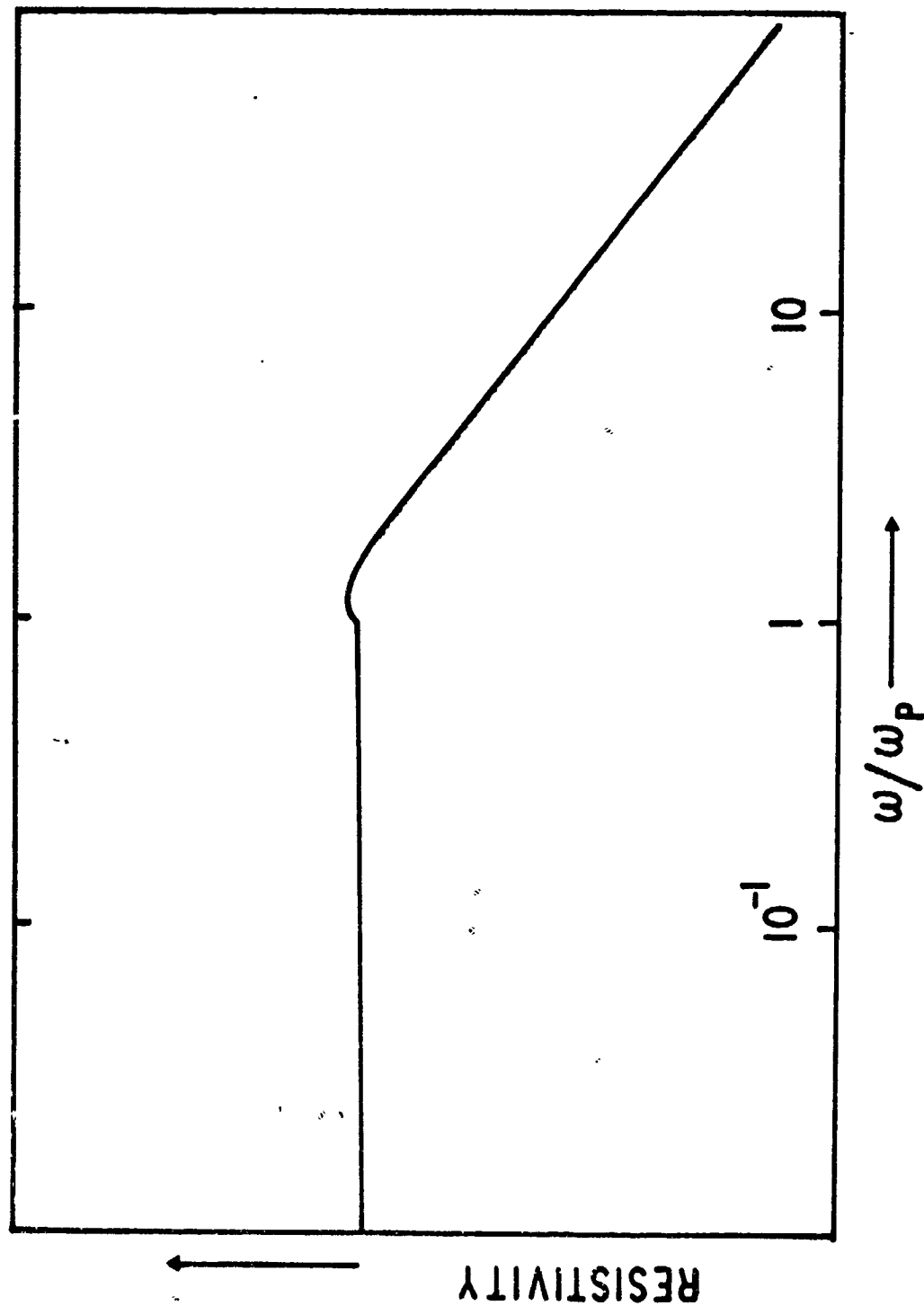
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Fig. 1. Dependence of the high-frequency resistivity on the normalized frequency (ω/ω_P).

AVERAGE OVER A PLASMA OSCILLATION

(a)

(b)
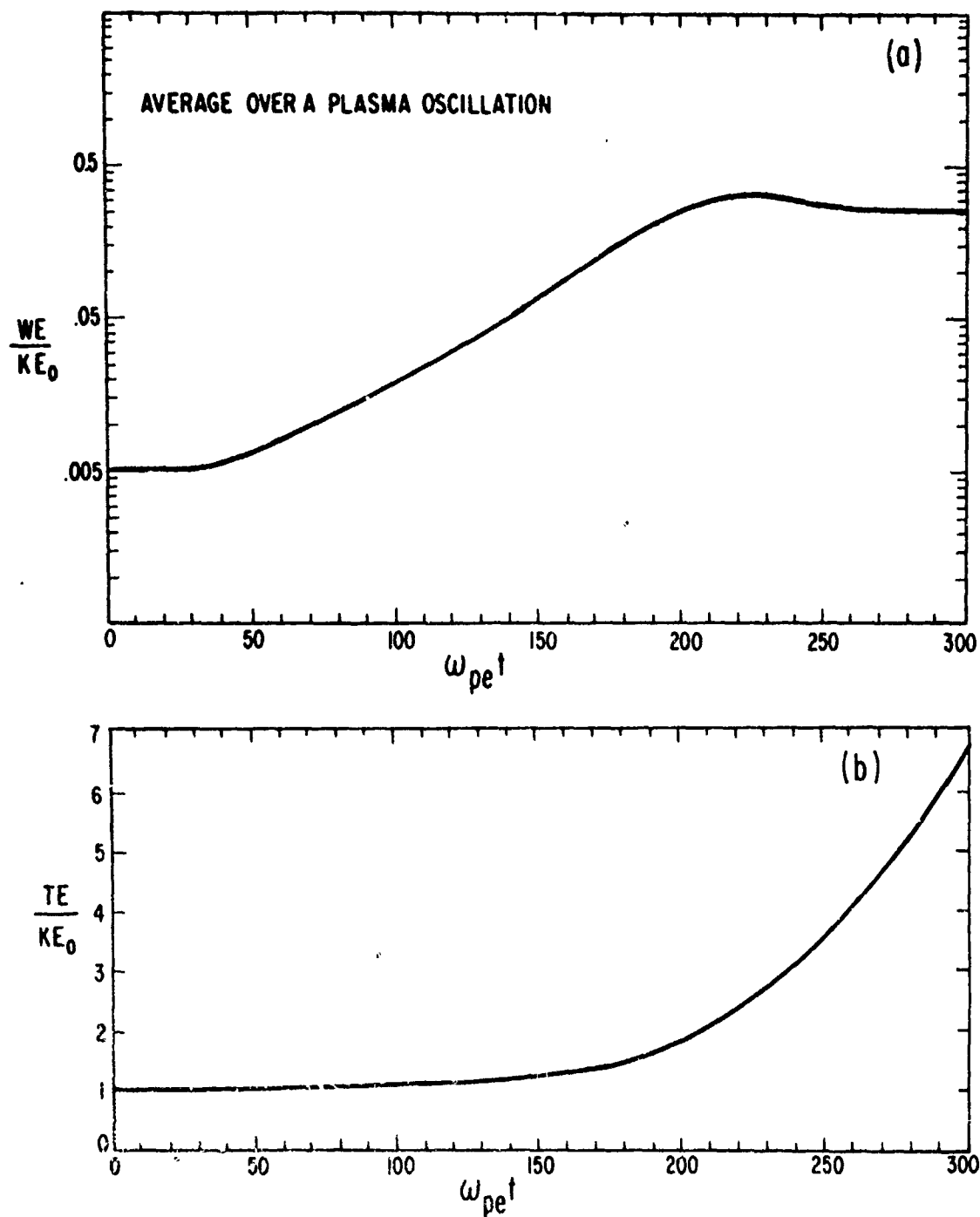
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Fig. 2. Results of the computer simulation experiment. The top trace is a plot of the total energy in the self-consistent fields normalized to the initial thermal energy vs $\omega_p t$. The lower trace is a plot of the normalized thermal energy vs $\omega_p t$.
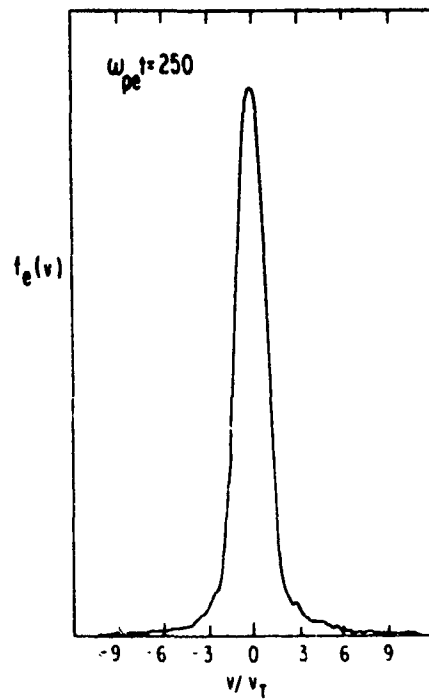
702117

Fig. 3. Electron velocity distribution after anomalous heating has taken place for $250\ \omega_p^{-1}$ sec.
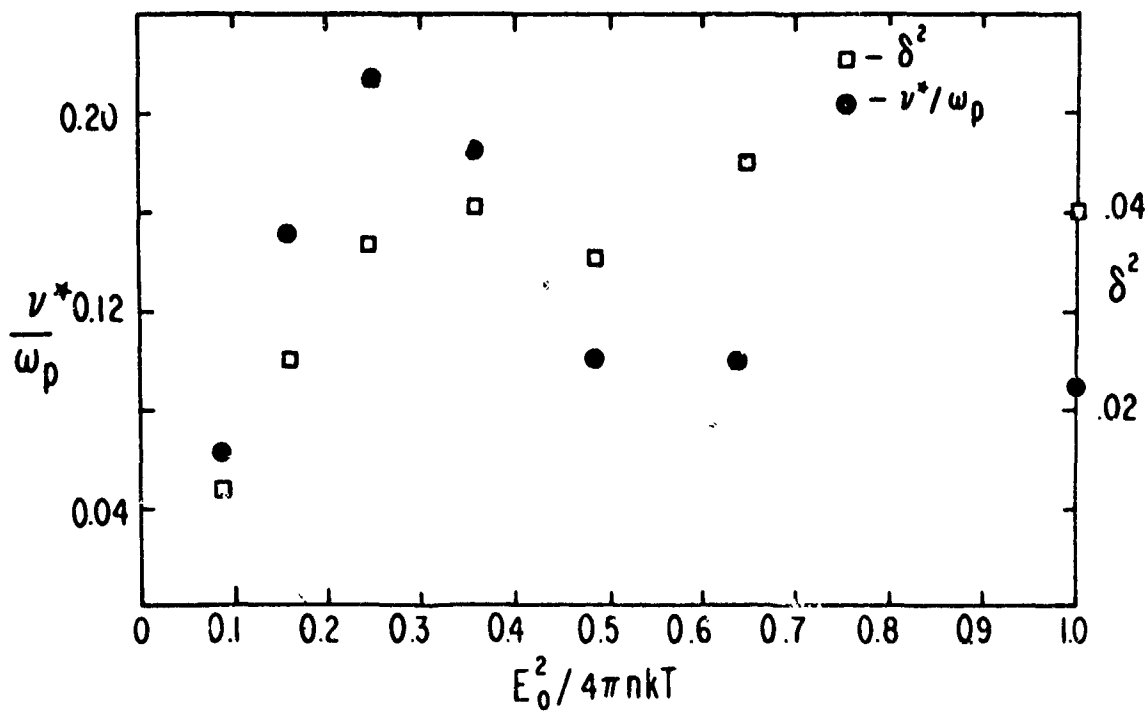
702023

Fig. 4. Variation of the normalized anomalous collision frequency (ν^*/ω_p) with the electric field parameter $(E_0^2/4\pi n\,KT)$ in the computer simulation experiment.
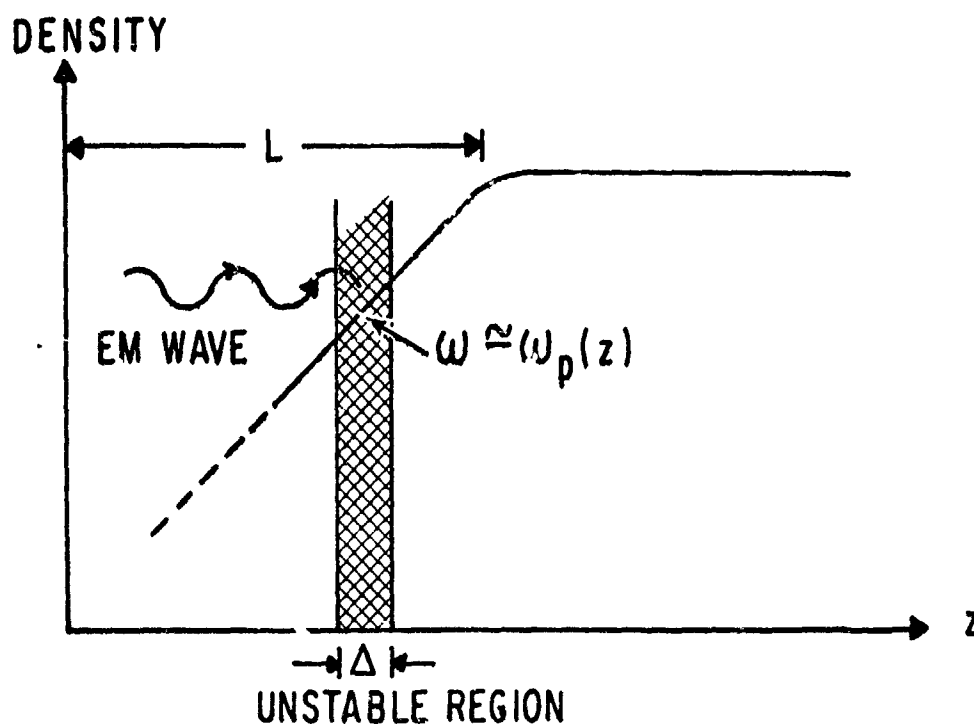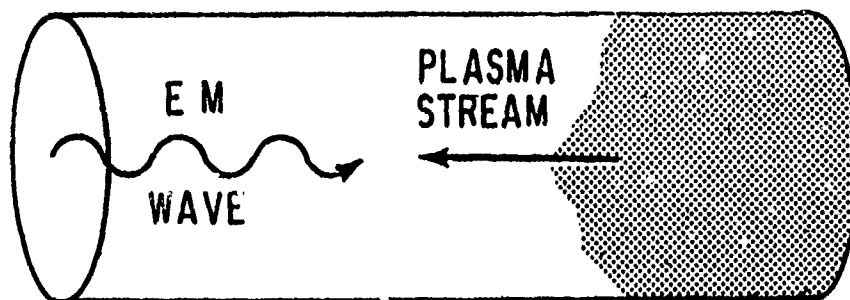
Fig. 5. Schematic of the Gekker-Sizukhin experiment.
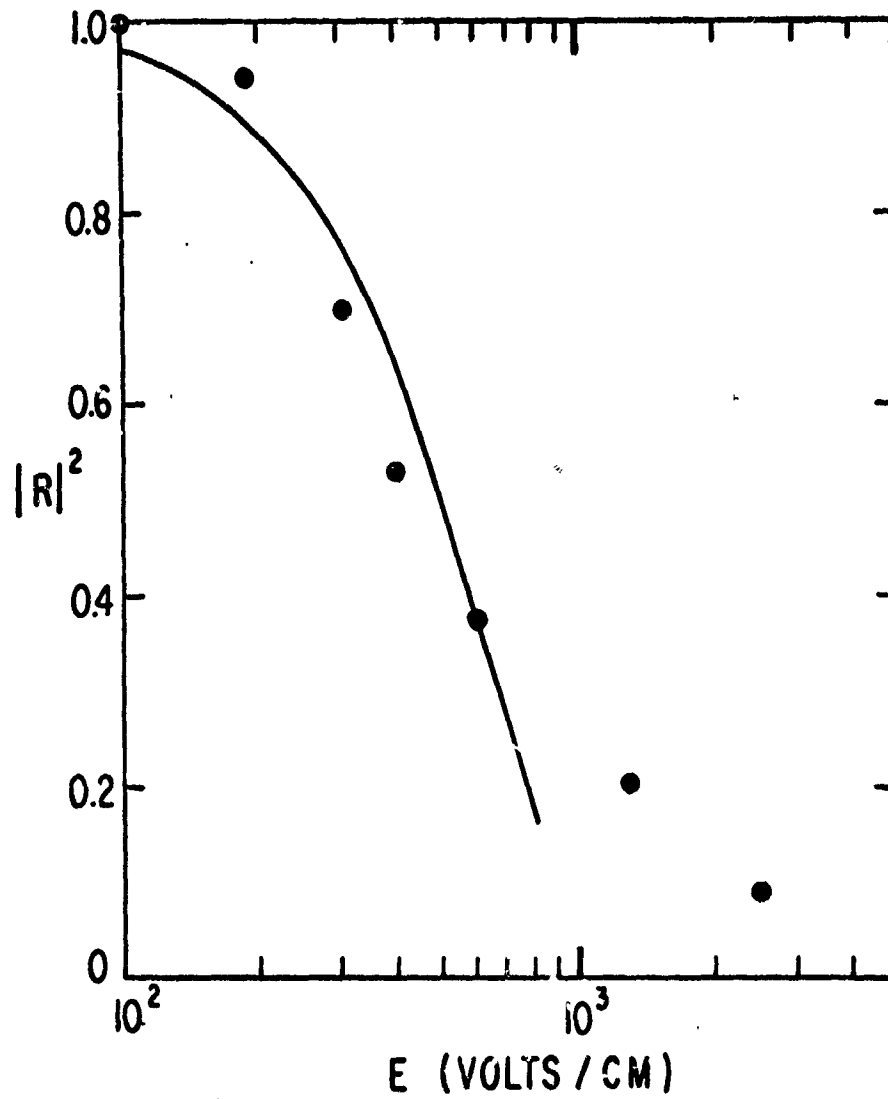
702098

702099

Fig. 6. Variation of the reflection coefficient $|R|^2$ with the microwave field amplitude E, from theory (solid curve) and experiment (black dots).

DENSITY PROFILE

ION PROBE

MICROWAVE
HORN

$\bigcirc$ B

$\vec{E} \parallel \vec{B}$
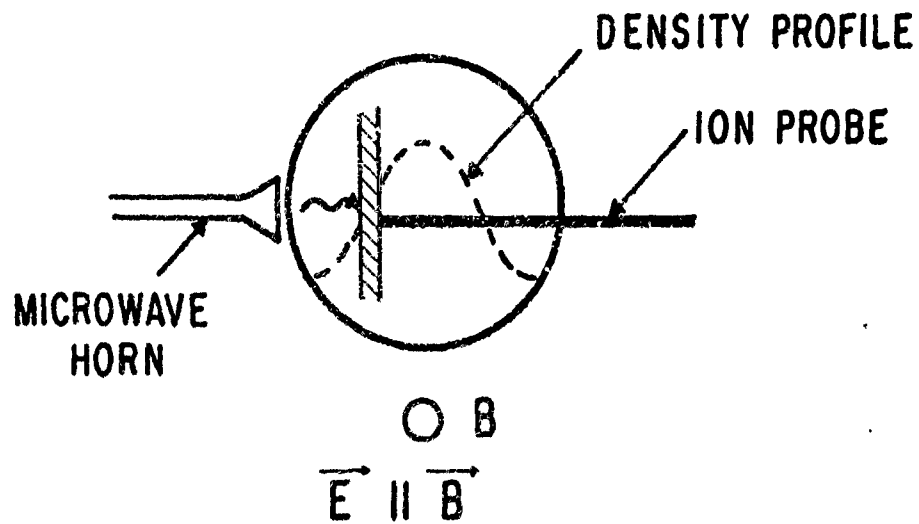
702093

Fig. 7.   Schematic of the Eubank experiment.

F − LAYER

T   POWERFUL   R
TRANSMITTER
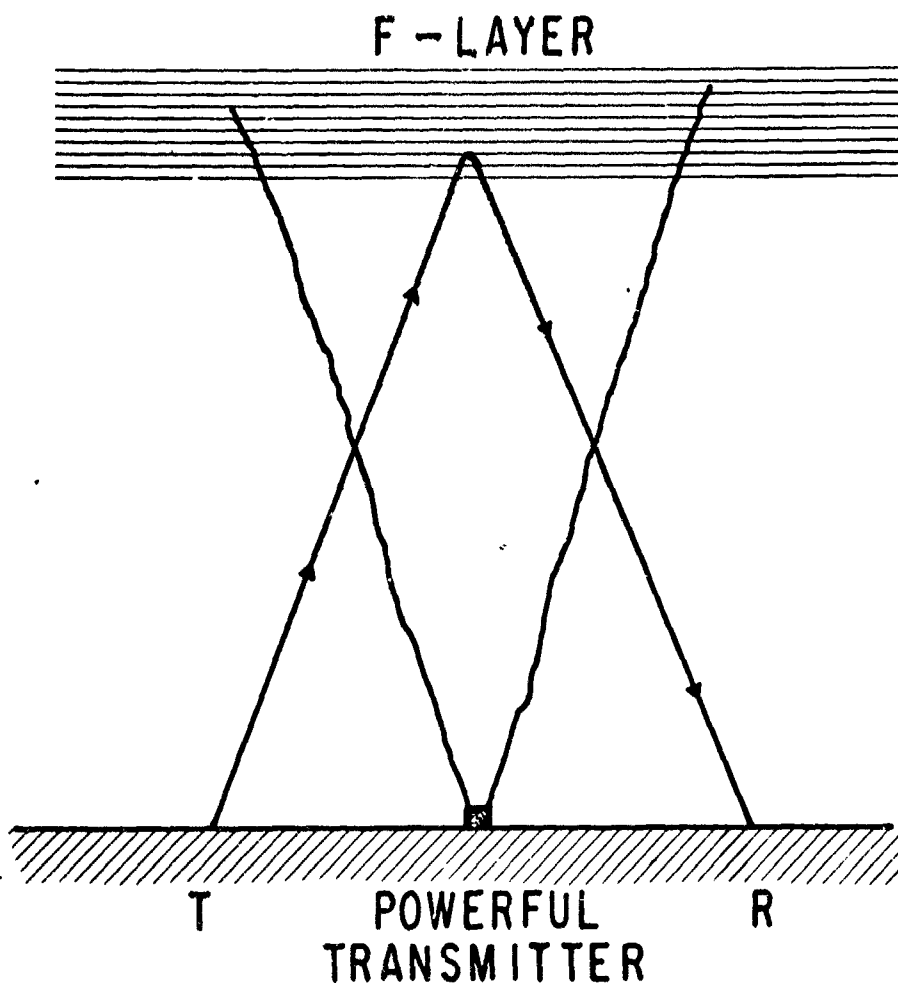
702091

Fig. 8.   Schematic of the ionospheric heating experiment.
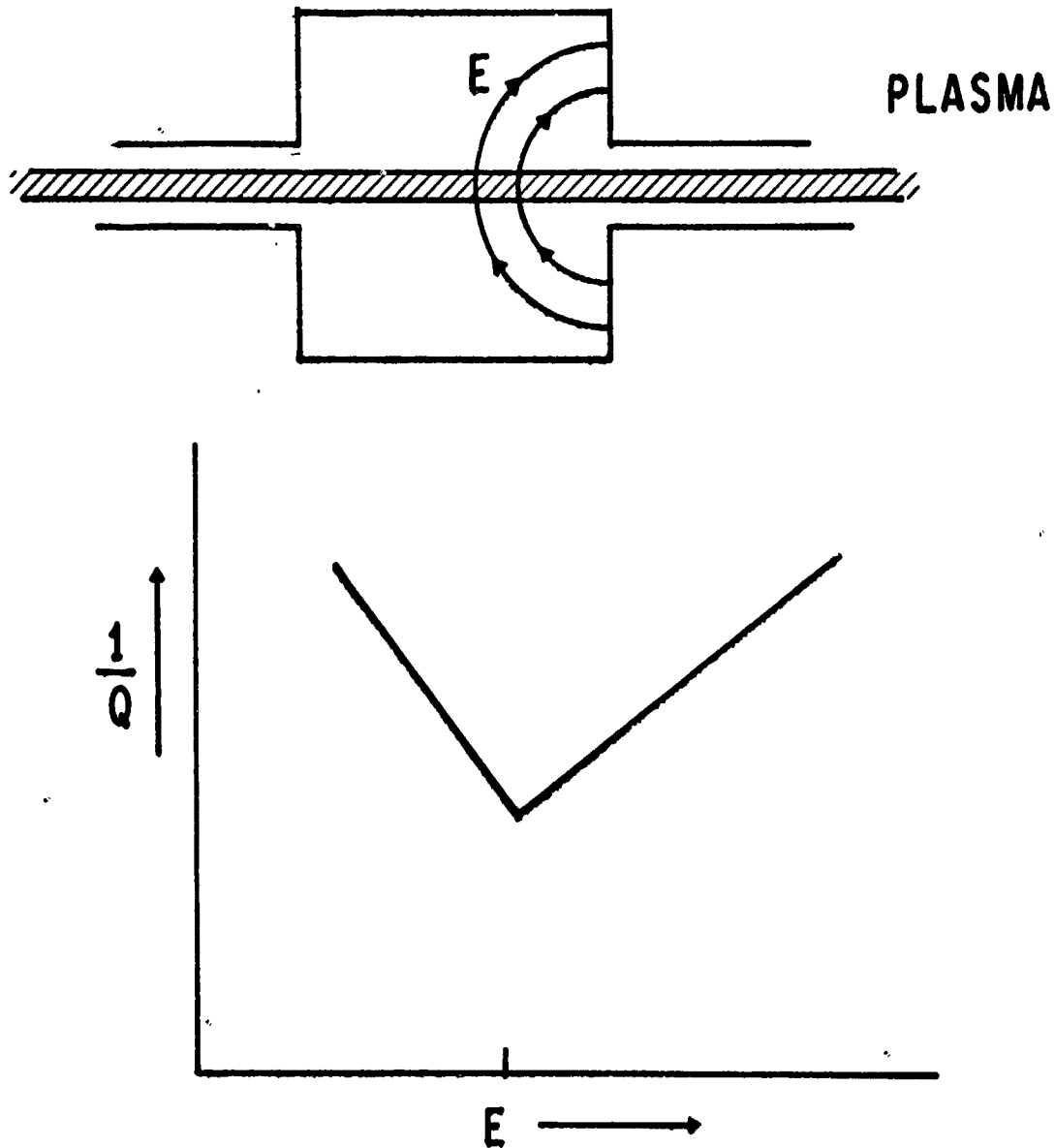
702094

Fig. 9. Schematic of the experiment due to Dreicer et al. The lower
trace is a plot of the dissipation (measured as the inverse of quality factor Q)
vs the microwave field amplitude.